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"INEVITABLE" DIBARYONS*

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Introduction

There exist many models of QCD: the MIT bag[1] and related variants such as the cloudy bag,[2] string and flux tube models,[3] non-relativistic quark cluster models with harmonic oscillator potentials,[4] and relativistic potential models such as the Los Alamos Model Potential[5] (LAMP). These models exist and proliferate because of the great difficulty of actually carrying out calculations in QCD, as Prof. 't Hooft discussed in his talk this morning.

The models differ greatly in technical details, but have three major features in common:

1. They implement color confinement, and so quark localization energies are an important component of spectral predictions.
2. They include spin dependence through the color magnetic spin-spin (CMSS) interaction (current-current in relativistic versions).
3. Parameters are fixed by fitting to the nucleon and delta, or more generally to octet and decuplet states.

Despite these commonalities, the predictions of these models also differ widely in spectral details: for higher resonances, for hybrid (quark-gluon) states, and in particular, for the H-particle of Jaffe.[6] This last ranges in predictions from deeply bound below $\Lambda\Lambda$ threshold[7] to so far above as to be unidentifiable due to a large width.

A point not often emphasized is that all predicted dibaryon spectra include an attractive effective interaction from the effect of quark delocalization. Even in cluster models, where delocalization is overtly excluded, a small amount is introduced perforce by Pauli (antisymmetrization) effects. Delocalization reduces the quark energies and contributes to binding of the state.

However, the widest variation in effects arises from the CMSS. In fact, since it tends to oppose quark delocalization by introducing new repulsive interactions (for example, between color-6 pairs of quarks), most dibaryon predictions are highly sensitive to the CMSS. What I want to discuss here are a few that are <u>not</u> sensitive, and, indeed, where the CMSS and quark delocalization act in concert to enhance binding.

The work reported here has been done in collaboration with G. J. Stephenson, Jr., K. E. Schmidt, K. Maltman, and Fan Wang.

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LLATH and CMSS

In what follows we restrict ourselves to six-quark, orbitally non-excited states. By the "lowest lying asymptotic two hadron" (LLATH) state of a given channel we will mean the lowest lying two baryon state, in a relative s-wave, having the given channel quantum numbers. The sum of the masses of the two baryons in the LLATH state then represents the threshold for fall- apart "decay" in the channel in question. A six-quark configuation lying below this threshold is necessarily resonant since it can be connected to a lower lying two-baryon state (by definition of the threshold, in a relative d- or higher wave) only by the action of some operator with non-trivial spin and orbital transformation properties. Moreover, since, as we see from the hadron spectrum, quark tensor forces are quite weak[8], such states will have small tensor decay widths[9] to any such lower lying two-baryon states and, in consequence, by very inelastic with respect such channels.

The color-spin structure of the hyperfine interaction is

$$-\left(\frac{3}{4}\right)\sum_{i<j}\vec{\sigma}_i\cdot\vec{\sigma}_j\ \ \lambda_i^a\lambda_j^a\ . \tag{1}$$

The structure of (1) is such as to favor color-spin symmetric pairs and hence, if we restrict ourselves to totally symmetric spatial configurations, lower flavor symmetries. The minimum value of the expectation of (1), as is well known, occurs in the H channel[6]. In the non-relativistic limit, the spatial dependence is a delta-function in the difference of quark coordinates. Relativistically, the effective range is short.

Let us now examine the expectation of (1) in the spatially symmetric and LLATH configurations for the channels we are considering. These values are given in Table 1, where $<>$ represents the expectation value of (1) and is to be multiplied by the corresponding spatial expectation in obtaining the total hyperfine energy.

Table 1: CMSS coefficients vs. channel

I	J	LLATH	$<>_{symm}$	$<>_{LLATH}$	Difference
0	1	NN	2	-12	14
1	0	NN	6	-12	18
1	2	ΔN	12	0	12
2	1	ΔN	20	0	20
3	0	ΔΔ	36	12	24
0	3	ΔΔ	12	12	0
$\frac{1}{2}$	2	NΩ	3	6	-3
$\frac{1}{2}$	1	ΛΞ	-7	-12	5
0	0	ΛΛ	18	-12	-6

One immediately sees that, for all channels but the last four cases, the hyperfine expectation is far more repulsive in the symmetric configuation than in the corresponding LLATII channel.

Furthermore, because of the short effective range of the CMSS interaction, the delocalization of the quarks in the 6-quark configuration reduces the CMSS spatial matrix element there, diluting its strength. This makes it clear that the $IJ^P = 03^+$ and $S = -3(J = 2)$ states actually experience an effective hyperfine attraction relative to the relevant two hadron thresholds. Since the dilution factor typically lies between 1 and 2, it is also clear that these are the only channels for which this is the case. (For the H, the net CMSS effect can still be repulsive accounting for some of the wide variation it its predicted mass.) Note, moreover, that <u>any</u> model incorporating a hyperfine interaction with the one gluon exchange color spin structure (1) and a delocalization mechanism in its dynamics, must necessarily produce a dibaryon resonance in these channels. For the effect of a dilution factor of 2, see Table 2.

Table 2: Effect of a factor 2 dilution (reduction of spacial matrix element) on CMSS contributions

I	J	LLATII	difference	net effect
0	3	$\Delta\Delta$	$[12(\frac{1}{2}) - 12] <>_3$	$-6 <>_3$
$\frac{1}{2}$	2	$N\Omega^-$	$[-3(\frac{1}{2}) - 0] <>_3$	$-1.5 <>_3$
$\frac{1}{2}$	1	$\Lambda\Xi$	$[-7(\frac{1}{2}) + 12] <>_3$	$+8.5 <>_3$
0	0	$\Lambda\Lambda$	$[-18(\frac{1}{2}) + 12] <>_3$	$+6 <>_3$

$<>_3$ represents the spatial matrix element in an isolated 3-quark state.

Calculations and Results

The LAMP[5] is a model which represents the confining structure of the QCD vacuum together with the effect of additional quarks and/or antiquarks in a hadron, by a linearly rising one-body Lorentz-scalar confining potential. Single quark wavefunctions are obtained by solving the Dirac equation for this potential. The underlying picture is bag-like, but with a gradual restoration of the full non-perturbative vacuum surrounding the hadron in question over the region in which the quark density becomes small. Since the structure of QCD is believed to be such that gluonic configurations characterizing the physical vacuum are suppressed in the presence of quark density, this means that, for dibaryons, an appropriate (mean field) potential can be obtained by truncation of the potentials from the individual baryon wells.

In the six quark sector,[10] rather than attempt to solve for the ground state of the truncated potential, we employ the following single-body trial wavefunctions

$$\psi_{\epsilon}(x) = [\psi(x - x_i) + \epsilon\psi(x - x_j)]/N(\epsilon) \tag{2}$$

where x_i is the center of the i^{th} potential well, $N(\epsilon)$ is a normalization factor, ψ is the $1S_{1/2}$ wavefunction in the single baryon well and, for $i = 1, 2, \quad j = 2, 1$. Although we restrict ourselves, here, to the case $\epsilon = 1$, such trial functions are, in general, a practical necessity, since the repulsive character of the color hyperfine interaction tends, in most channels, to produce localized quark structure, the wavefunctions of which may be represented only by a superposition of many single-particle levels of the truncated, two-centered potential. Note, however, that because $\epsilon = 1$, these states are not in any sense "nuclear," i.e. composed of baryons. They represent a distinctly new structure of hadronic matter.

The results of our calculations are shown in Table 3, where d is the separation between the centers of the potentials for the individual baryons for which the minimum of the total energy (single body plus CMSS) occurs. For the smaller values of d, our ansatze may well be called into question, as the quark density in the interbaryon region can become rather large. This suggests the potential should be further suppressed, but that should in turn add to the overall energy of the system, since "perturbative QCD vacuum" is being restored. This raises questions about overestimating the binding energies. Nonetheless, the characteristic "peanut" or "football" shapes in these body fixed frames (the individual baryon rms sizes are ~0.8 fm) seem likely to survive more detailed analysis.

Table 3: Binding Energies and Size Scales

State	Binding Energy (MeV)	LLATH (threshold)	d (fm)
Deltaron ($J = 3$)*	~350	$\Delta\Delta$	1.4
Omegon ($S = -3$, $J = 2$)†	290±10	$N\Omega$	0.7 to 0.8
Omegon ($S = -3$, $J = 1$)†	190±30	$\Lambda\Xi$	0.7 to 0.8
H($S = -2$, $J = 0$)†	210±110	$\Lambda\Lambda$	0.5

* See Ref. 10. † See first paper in Ref. 5.

In Table 4, we present the results of deltaron†† calculations in other models of confinement. Obviously deviation from sphericity enhances binding, as for the deuteron itself. In view of the concordance of these results, we find it highly plausible that the d^* occurs very close to

Table 4: Deltaron in Other Models of Confinement

Model	Binding Energy (MeV)	Comments
MIT bag[1]	115	restricted to spherical state
Cloudy bag[2]	85	restricted to spherical state
Non-relativistic quark model[4]	290	non-spherical but no quark delocalization, center of mass motion removed

††We name it so because its quantum numbers are those of a spin excitation of the deuteron and denote it by d^*.

$NN\pi\pi$ threshold. The LAMP would even have it slightly below. In either event, the width is likely to be small — to $NN\pi\pi$ because of the small phase space and to $NN\pi$ because a quark spin flip is still required. We have made several estimates of the deltaron width, including off-shell Δ-decays; these all turn out to be less than a few 10's of MeV.

Conclusions

I have had neither time nor space to fully discuss appropriate caveats regarding the particular calculations. We take the excellent results for the LAMP applied to ^{4}He (see second paper of Ref. 5) as some support for our calculations. Futhermore, both the universal nature of the deltaron prediction and our understanding of the reason for this (from the clarity and flexibility afforded in the LAMP calculation) give us confidence that the state exists. It and the $J = 2$ omegon should be obervable. The latter has a long enough (predicted) lifetime to appear in hyperon beams. While not easy to see in the NN $^3D_3$ wave due to the smallness of quark tensor forces, the deltaron should be observable in processes such as $\pi d \rightarrow \pi d^*$, $ed \rightarrow e'd^*$ or $e'd^*\pi$ and $pd \rightarrow pd^*$ or $Nd^*\pi$.

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